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THERMAL-CHEMICAL INSTABILITIES IN CO CLOUDS

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ABSTRACT

The stability of interstellar clouds containing CO is investigated taking into account formation and destruction processes for molecules. Thermal-chemical instabilities are obtained which influence the evolution of clouds.

Subject headings: interstellar: matter — interstellar: molecules — molecular processes

Carbon monoxide is widely regarded as an important source of cooling in interstellar clouds (Rank *et al.* 1971). Earlier studies have been concerned with the effectiveness of CO cooling at constant abundance in a variety of situations (Glassgold and Langer 1973; Berger and Simon 1973; Scoville and Solomon 1974; Goldreich and Kwan 1974; de Jong *et al.* 1975). When formation and destruction processes for CO are taken into account, however, new aspects of the thermal and dynamic properties of the clouds may be expected to occur. In this report we investigate some of these effects, specifically the question of whether the conversion from C^+ to CO leads to instabilities which influence the evolution of clouds, as suggested by Glassgold and Langer (1975) and by Oppenheimer and Dalgarno (1975).

In order to study the stability of the material in molecular clouds, a generalization of Field's paper (1965) which includes chemical effects must be used. We will only discuss a few essential aspects of this generalization here; further details can be found in Goldsmith (1970), Defouw (1970), and Yoneyama (1973) who have discussed related problems. The basis of the stability analysis is the linearized form of the conservation laws of hydrodynamics, augmented by equations of state. The equations for mass and momentum conservation and the ideal gas law are the same as Field's equations (11), (12), and (14), because the heavy atoms and the molecules formed from them have negligible density compared with the bulk of the gas. The energy equation, $du/dt = p\rho^{-2}d\rho/dt - \mathscr{L}$ must be generalized to include chemical contributions to the internal energy per unit mass u and to the rate of net energy loss per unit mass $\mathscr{L}$; $\mathscr{L}$ must also contain cooling from C^+ and CO as well as appropriate heating mechanisms. Finally, there must be a rate equation $dx_i/dt = J_i$ specifying the rate of change of x_i, the abundance relative to hydrogen of each chemical species. However, we have been able to show that, for the chemical schemes considered here, the time scales for all species are considerably shorter than those for C^+ and CO, except for H_2 which has a much longer one (Langer and Glassgold 1976). Consequently, only the time dependence of CO needs to be explicitly solved for in first approximation. An exception to this can arise at low temperatures and low densities where a considerable amount of carbon is in C I, which is not the case for the situations discussed here. Thus the required extension of the instability theory involves a single additional hydrodynamic equation for $x \equiv x(\mathrm{CO})$,

$$dx/dt = J(T, \rho, x)\,. \tag{1}$$

The abundance of H_2 can be considered as constant, whereas the abundances of other species such as OH^+, OH, $CH_2{}^+$, C_2H, etc., are obtained from steady-state formulae in which the CO abundance x is a variable. Because most of the recombining C^+ goes into CO, $\dot{x}(C^+) = -\dot{x}$ and $x(C^+) \approx \xi_C - x$, where ξ_C is the abundance of gaseous carbon.

The characteristic equation for the growth of perturbations in a uniform medium without heat conduction or viscous dissipation can be obtained by substituting solutions proportional to exp $(i\vec{k}\cdot\vec{r} + \sigma t)$ into the linearized system of hydrodynamic equations and equations of state. The result in this case is the quartic equation

$$\sigma^2[\sigma^2 + (\kappa_\rho + \mu_0)\sigma + \mu_0\lambda_\rho] + c^2k^2[\sigma^2 + (\kappa_p + \mu_0)\sigma + \mu_0\lambda_p] = 0\,, \tag{2}$$

where c is the adiabatic sound speed, κ_i and λ_i are inverse thermal time scales, defined by

$$\lambda_i = \frac{1}{c_i}\left[\left(\frac{\partial\mathscr{L}}{\partial T}\right)_{i,x} + \mathscr{L}_x\left(\frac{\partial x_0}{\partial T}\right)_i\right],$$

$$\kappa_i = \frac{1}{c_i}\left[\left(\frac{\partial\mathscr{L}}{\partial T}\right)_{i,x} - \mu_0\Delta E\left(\frac{\partial x_0}{\partial T}\right)_i\right] \quad (i = p, \rho)\,, \tag{3}$$

and $\mu_0 > 0$ is the inverse chemical time scale obtained from the linearized form of equation (1),

$$\frac{dx}{dt} = -\mu(T, \rho, x)(x - x_0) \approx -\mu_0(x - x_0) . \quad (4)$$

In equation (3), x_0 is the steady-state abundance, c_i is the heat capacity per unit mass at constant $i = p$ or ρ, $\mathscr{L}_x = (\partial\mathscr{L}/\partial x)_{T,\rho}$, and ΔE is the energy difference per unit mass between C^+ and CO (proportional to the sum of the ionization potential of C I and the dissociation energy of CO). When chemical effects are unimportant, $\mu_0 \to 0$ and equation (2) reduces to Field's characteristic equation without heat conduction. More generally, a sufficient condition for instability is a negative constant term in the quartic, i.e.,

$$\lambda_p = \frac{1}{c_p}\left[\left(\frac{\partial\mathscr{L}}{\partial T}\right)_{p,x} + \mathscr{L}_x\left(\frac{\partial x_0}{\partial T}\right)_p\right] < 0 . \quad (5)$$

The quartic equation (2) has been written in terms of two quadratic forms $d_i = \sigma^2 + (\kappa_i + \mu_0)\sigma + \mu_0\lambda_i$, $i = p, \rho$. The characteristic equation reduces to $d_p(\sigma) = 0$ for wavelengths sufficiently short that the perturbation becomes isobaric within the time interval $|\sigma|^{-1}$, i.e., for $|\sigma| \ll ck$; for $|\sigma| \gg ck$, the characteristic equation is $d_\rho(\sigma) = 0$.

We will apply these instability considerations to the material in relatively diffuse clouds with hydrogen gas densities $n = n(H) + 2n(H_2)$ in the range from $\sim 10^2$–10^3 cm^{-3} and cloud thicknesses such that $A_v \lesssim 3$. For a slab model, the column density N and distance L to the center are in the range $N \lesssim 3.75 \times 10^{21}$ cm^{-2} and $L < 12.5$ pc. The results will depend on the type of chemistry assumed, and we have investigated two gas-phase ion-molecule reaction schemes:

1. O^+ *charge exchange*. CO is formed by reactions of C^+ with OH and H_2O, the latter being produced by a complex of reactions of O^+ with the weakly ionized H_2 gas. Because the O^+ ions are formed by the slightly endothermic charge-exchange of cosmic-ray produced H^+ ions with O, this formation mechanism becomes very weak at low temperatures (below $\lesssim 40$ K). This scheme was proposed by Watson (1973, 1974*a*, *b*) and Dalgarno and collaborators (Dalgarno *et al.* 1973; Black and Dalgarno 1973*a*), and further developed by Glassgold and Langer (1976) and by Oppenheimer and Dalgarno (1975). As long as CO does not exhaust all of the available carbon, the chemical time scale for the situation of interest here is roughly given by (Langer and Glassgold 1976)

$$\mu_0 \approx n\gamma \equiv G(CO) + k_{25}x(He^+)n , \quad (6)$$

where $G(CO)$ is the photodissociation rate of CO and k_{25} is the rate constant for destruction of CO by ion-molecule reactions with He^+.

2. C^+ *radiative association*. CO is formed by neutral molecule reactions of O with CH, CH_2, C_2H, and C_2, which are produced by a complex of reactions initiated by C^+ with the weakly ionized H_2 gas. It is assumed that this process is not inhibited at very low temperatures. This scheme was proposed by Black and Dalgarno (1973*b*), and further developed by Oppenheimer and Dalgarno (1975) and by Langer (1976). The chemical time scale in this case is roughly given by (Langer and Glassgold 1976)

$$\mu_0 \approx [\tfrac{1}{2}k_{26} + k_{25}x(He^+)]n \quad (7)$$

where k_{26} is the rate constant for radiative association of C^+ and H_2.

The following form has been used for the energy loss:

$$\mathscr{L} = \mathscr{L}_{\rm cool} + \mathscr{L}_{\rm Gr} + \mathscr{L}_{\rm CR} + \mathscr{L}_{\rm ch} . \quad (8)$$

$\mathscr{L}_{\rm Gr}$ and $\mathscr{L}_{\rm CR}$ are heating terms associated, respectively, with photoelectrons ejected from grains by the interstellar ultraviolet radiation field and with cosmic-ray protons interacting with the hydrogen-helium gas mixture. The specific forms used here were discussed earlier (Glassgold and Langer 1974). The grain term is assumed to vary roughly as $\exp(-\tau_{\rm Gr})$, where $\tau_{\rm Gr}$ is the optical depth for absorption by grains near 1100 Å. We use a primary cosmic-ray ionization rate for H of $\zeta_p = 5 \times 10^{-17}$ s^{-1} unless otherwise noted. $\mathscr{L}_{\rm ch}$ is the rate per unit mass at which the energy changes due to chemical reactions, and has been called chemical heating by Oppenheimer and Dalgarno (1974). We have not found $\mathscr{L}_{\rm ch}$ to be very important in the present context and, for purposes of brevity, omit further discussion of this term. $\mathscr{L}_{\rm cool}$ is the collisional-radiative cooling rate due to C^+ fine structure and CO rotational excitations.

In addition to the familiar excitation of the fine-structure doublet by H atoms and electrons, we also include collisions with H_2 molecules (Chu and Dalgarno 1975). For the CO cooling rate per unit volume, we use (Leung 1975)

$$L = \sum_{J'=1}^{\infty}\sum_{J=0}^{J'-1} E_{J'J}(n_J C_{JJ'} - n_{J'}C_{J'J}) , \quad (9)$$

where n_J is the density of CO molecules with rotational quantum number J, $E_{J'J}$ is the energy difference between two such levels, and the C's are the transition rates due to collisions with H_2 molecules. The n_J are the steady-state populations obtained by considering simultaneously these collisions and radiative transitions, including the 3 K background radiation field. This point of view is similar to that adopted by de Jong *et al.* (1975) and by Oppenheimer and Dalgarno (1975), except that we consider a stationary medium. We have used the most recent calculations for the $C_{JJ'}$ due to Green (1975) in which $\Delta J \leq 9$. The effect of trapping of the cooling radiations can be important for some of the larger optical depths considered here and will be discussed below.

We begin our discussion of results with the limiting case $|\sigma| \ll ck$ for which a thermodynamic analysis is appropriate. The sign of the slope of a pressure versus density curve is the same as the sign of λ_p in the instability condition equation (5) (Field 1975). Thus an instability is indicated by a loop in a pressure

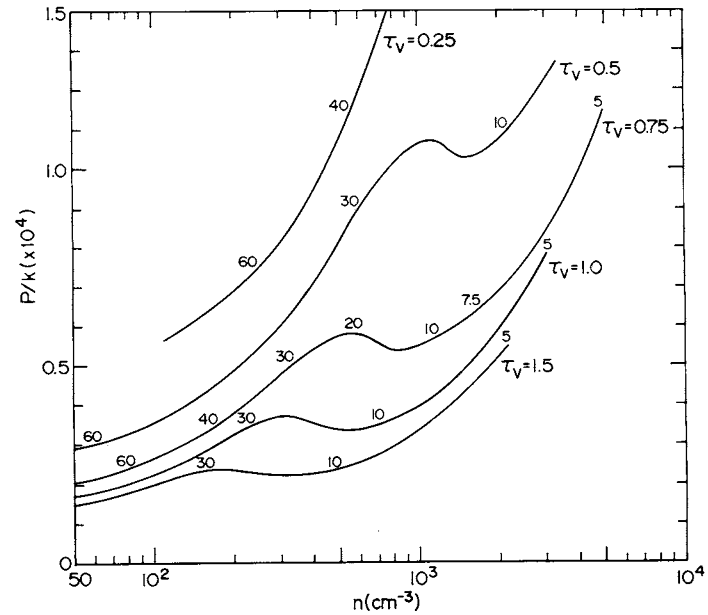

FIG. 1.—The pressure, $P/k = n(1 - f/2)T$, is plotted as a function of hydrogen density for a range of optical depths into a cloud, τ_v. The parameters chosen for this set of curves are the following: $\xi_O = 2 \times 10^{-4}$, $\xi_C = 7 \times 10^{-5}$, $\zeta_p = 5 \times 10^{-17}\ s^{-1}$, $k_{26} = 10^{-15}\ cm^3\ s^{-1}$, $G_{CM}(\tau_v = 0) = 10^{-10}\ s^{-1}$, $G_{CO}(\tau_v = 0) = 3 \times 10^{-12}\ s^{-1}$, and $f = 0.98$. The numbers along each curve label the temperature.

curve. Examples of such curves for the CO chemistry discussed earlier are given in Figure 1 for several optical depths with all other parameters fixed; $\tau_v = A_v/2 = N/(2.75 \times 10^{21}\ cm^{-2})$, where τ_v and N are to the center of a cloud. Instabilities are indicated at low temperature for $T \approx 20$ K for $\tau_v \gtrsim 0.5$. Even for $\tau_v \approx 0.25$ an extension of the curve in the figure would show a very soft equation of state at $T \approx 20$ K, i.e., a region where λ_p is positive but very small.

The qualitative features of Figure 1 are relatively insensitive to the choice of parameters, as will be discussed below. Some of the more important parameters for this figure are the following: $\xi_C = 7 \times 10^{-5}$ and $\xi_O = 2 \times 10^{-4}$ (abundances), $f \cong 1.0$ ($f = 2n(H_2)/n$), and $k_{26} = 10^{-15}\ cm^3\ s^{-1}$. This value of f is appropriate for $\tau_v > 0.5$ and $n \gtrsim 100\ cm^{-3}$, and its choice influences C^+ cooling, because atomic hydrogen has a larger cross section for exciting the fine-structure transition than does H_2. The way in which k_{26} occurs in the low-temperature chemistry may be seen from an approximate form of equation (1)

$$J \approx n\left[\frac{f}{2} k_{26}(\xi_C - x)b - \gamma x\right], \qquad (10)$$

where $b = ax(O)/[G_{CM}/n + ax(O)]$ is a branching ratio determined by a characteristic photodissociation rate for the CH family of molecules and by a, the rate constant for forming CO from these molecules in reactions with O [abundance $x(O)$]. For small optical depths the effectiveness of the radiative association mechanism for forming CO is measured by the combination $k_{26}na/(G_{CM}\gamma)$. Although all of these parameters are uncertain, variations in optical depth and density generally lead to regions with substantial CO abundance.

The main features of Figure 1 can be qualitatively understood in terms of the temperature and density dependence of C^+ and CO cooling and the density dependence of the CO chemistry at low temperatures. When T is reduced toward 30 K, C^+ cooling decreases sharply because $\mathscr{L}_{cool}(C^+) \propto ne^{-91.2/T}$. The density n will undergo a corresponding increase along one of the pressure curves in Figure 1 in order to maintain thermal balance, the heating per unit mass being roughly constant. If there is little CO production, the equation of state is then very stiff. This behavior can be seen clearly in the curves for small τ_v as T ranges from 60 to 30 K. A limiting minimum temperature would eventually be reached for each pressure curve (or value of τ_v), because collisional de-excitation will make $\mathscr{L}_{cool}(C^+)$ independent of n for sufficiently large n. A similar situation occurs at much lower temperatures $\sim$5 K where the cooling is dominated by only one or two of the lowest levels of CO, and $\mathscr{L}_{cool}(CO) \propto n \exp(-T_c/T)$ with $T_c \approx 6$ K. In other words, the equation of state also becomes stiff at very low temperatures, and a limiting temperature $T_\infty \approx 4$ K would eventually be reached because of collisional de-excitation (Glassgold and Langer 1973). Trapping of the cooling radiation from the low-lying levels could raise the minimum value T_∞. Other molecular coolants or gas cooling by collisions with cold dust particles might also be important at the high densities required for the achievement of very low temperatures below 10 K.

Below about 30 K, temperature reductions are accompanied by relatively large increases in density.

Because the cooling efficiency (cooling rate per carbon nucleus) is greater for CO than for C^+, an instability may occur if CO production increases sufficiently with density. The condition for the onset of strong CO production is, from equation (10), $b \approx \frac{1}{2}$, which determines a characteristic value of the density $n' = G_{CM}/ax(O)$. The shift of the instabilities toward smaller n as τ_v is increased in Figure 1 arises from the decrease of G_{CM} with τ_v. Similarly, if somewhat different parameters are used instead of those selected for Figure 1, the instabilities will still occur but at shifted locations. For small optical depths, n' becomes so large that the CO cooling efficiency is significantly reduced by collisional de-excitation, with the consequence that the thermal anomaly is no longer an instability but simply a very soft equation of state.

The above qualitative discussion of the loops in the pressure curves in Figure 1 is borne out by a quantitative study of the roots of the characteristic equation (2). In particular, the instability condition in equation (5) is a statement of how strong the thermal effect of converting C^+ to CO has to be to produce an instability, i.e., $\mathscr{L}_x(\partial x/\partial T)_p < -(\partial \mathscr{L}/\partial T)_{p,x}$. The radiative-association chemistry leads to low temperature instabilities in Figure 1, because $(\partial x/\partial n)_T$ dominates $(\partial x/\partial T)_p = (\partial x/\partial T)\rho - nT^{-1}(\partial x/\partial n)_T$. For the case $\tau_v = 0.75$ in the limit $|\sigma| \ll ck$, the following numerical results have been obtained on the roots of equation (2) for thermal-chemical modes as distinct from ordinary sound waves. At 40 K, there are two stable solutions with $|\sigma|^{-1} = 2 \times 10^5$ and 4×10^4 yr, and at 5 K the two stable modes have $|\sigma|^{-1} = 3 \times 10^4$ and 5×10^3 yr. At $T = 20$ K, there is one stable mode with $|\sigma|^{-1} = 7 \times 10^4$ yr and one unstable mode with $|\sigma|^{-1} = 4 \times 10^5$ yr. For smaller values of k, we find for the unstable mode that $|\sigma(k)|^{-1}$ varies from $\sim 9 \times 10^5$ to 7×10^6 yr, as k^{-1} ranges from 0.1 to 1 pc, which corresponds to a range of masses from $\sim$1–1000 $M_\odot$.

Trapping of the C^+ and CO line radiations decreases the cooling rate and thus may modify the regions where instabilities occur. Although a detailed treatment of trapping is beyond the scope of this paper, a qualitative discussion can easily be given. The trapping of the C^+ cooling radiation is negligible here for two reasons: (1) the opacity at 157 μ is small ($\lesssim \frac{1}{2}$), and few of these photons are absorbed; (2) the local densities considered here $n < 2 \times 10^3$ cm^{-3} are at least an order of magnitude less than the critical density for collisional de-excitation, $n_c \approx 2 \times 10^4$ cm^{-3}. Thus the probability for an absorbed photon to be subsequently reemitted is very large.

The optical depth for line absorption by a column density N(CO) of CO is

$$\tau(J \rightarrow J+1)$$
$$\approx 3 \times 10^{-16} N(\mathrm{CO}) \left(\frac{J+1}{2J+1}\right) P_J (b_D/3\ \mathrm{km\ s^{-1}})^{-1}, \quad (11)$$

where P_J is the fractional population of the Jth level and b_D is the usual parameter specifying the Doppler width. Using integrated column densities (Langer 1976) from the regions of instability to the cloud exterior, we find: $\tau(0 \rightarrow 1) \approx O(10)$ (optically thick); $\tau(1 \rightarrow 2) \approx O(1)$ (some absorption); $\tau(J \rightarrow J+1) \ll 1$ for $J \geq 2$ (optically thin). Again, whether an absorbed $0 \rightarrow 1$ or $1 \rightarrow 2$ photon subsequently escapes to cool depends on the critical density for collisional de-excitation through an efficiency factor for each level of the form $[1 + k_J(T)n_c(J)/n]^{-1}$ where $k_J(T) \approx O(1)$. The critical density is $n_c(J) = \beta_J n_c^{(0)}(J)$ where β_J is the escape probability and $n_c^{(0)}(J) = A_{J,J-1}/C_J$ is the critical density when there is no trapping (the ratio of the spontaneous transition probability to the total collisional rate out of the level); e.g., $n_c^{(0)}(1) \approx 2 \times 10^3$ cm^{-3} and $n_c^{(0)}(2) \approx 8 \times 10^3$ cm^{-3}. These considerations lead to the conclusion that, for the optical depths and densities of interest, cooling from the $1 \rightarrow 0$ transition is strongly suppressed, but cooling from the $2 \rightarrow 1$ transition is reduced slightly, and by no more than 0.5 for small τ_v. The thermalization of the $1 \rightarrow 0$ transition alters the population of the $J = 1$ level so that there is increased cooling from $J = 1 \rightarrow J' > 1$ transitions which tends to compensate for the $2 \rightarrow 1$ trapping. A reasonable approximation to the trapping in this problem involves removal of some or all of the $1 \rightarrow 0$ cooling. When this is done, the essential features of Figure 1 remained unchanged. The curves do become somewhat flatter, and the range of τ_v for which instabilities are present is reduced slightly. In other words, most of the cooling given by equation (9) without trapping arises from downward transitions with $J' \geq 2$ which are negligibly trapped for the situations of interest in this paper.

The warm low-density states ($n \lesssim 100$ cm^{-3}, $T \lesssim 100$ K) often associated with diffuse or "standard" interstellar clouds are to be found in the lower left portion of a p versus n phase diagram. The extrapolation of the curves of Figure 1 into this region is somewhat of an oversimplification, because they were obtained using a fixed value of f, the H_2 fraction. Pressure curves in this region do not manifest thermal anomalies in the usual sense. A study of the instability condition equation (5) using the O^+ charge exchange chemistry leads, however, to another type of instability associated with attenuation of the interstellar radiation field by grains. This radiation enters most importantly here through grain photoelectron heating and photodissociation of molecules. The rates of these processes decrease going into a cloud, giving rise to derivatives with respect to density which affect the characteristic equation. When a density fluctuation occurs, it will be accompanied by a change in optical depth. For example, if the CO photodissociation rate is approximated by $G(\mathrm{CO}) = G_0(\mathrm{CO}) \exp[-\tau_{Gr}(\mathrm{CO})]$, its density derivative is proportional to $[\tau_{Gr}(\mathrm{CO})/n] \times G(\mathrm{CO})$. This effect cannot be completely correctly dealt with by a uniform-medium instability analysis, which is only appropriate in the limit that the linear size of the perturbation is small compared with an attenuation length, i.e., for wavelengths $\lambda \ll \kappa_{Gr}^{-1}$, where κ_{Gr} is an appropriate absorption coefficient; for $n = 100$ cm^{-3}, $\kappa_{Gr}^{-1} \approx 3$ pc. When this effect is included in the instability analysis, both terms in equation (5) can become negative and produce an instability. The new contribution to the first term arises from

$$(\partial \mathscr{L}_{\mathrm{Gr}}/\partial T)_p = -n/T(\partial \mathscr{L}_{\mathrm{Gr}}/\partial n)_{T,x} < 0 ;$$

similarly, the second term in equation (5) becomes negative through the dominance of $(\partial x_0/\partial n)_{T,x}$ associated with the variation of the CO abundance with optical depth. A numerical study for $n = 100$ cm^{-3} indicates one unstable mode for temperatures in the range 40–80 K and for $1 \leqslant A_v \leqslant 3$, with a time scale typically $\sim 2 \times 10^5$ yr for smaller perturbations with $\lambda < 0.1$ pc. The other chemical mode is stable with a time scale 3 or 4 times shorter. Finally, it can be shown that this type of instability vanishes for small optical depths. Diffuse clouds with small enough A_v will be stable with respect to the thermal-chemical mechanisms discussed here.

The optical depth instability is not associated with any single curve in Figure 1, but with the changes between curves for different optical depths. Furthermore, the curves in Figure 1 suggest that this effect is not restricted to high temperatures, and a detailed study bears this out. When these effects are included in the analysis of the loops in the pressure curve, the instabilities are enhanced. However, there is always a low temperature, $\sim$5–6 K, at which this enhancement vanishes and the curves are stable again. In other words, this thermal-chemical instability vanishes for large as well as small optical depths.

Thermal-chemical instabilities can also occur if grain formation rather than ion-molecule chemistry dominates CO production. In this case a model similar to that used for the radiative association chemistry is applicable and $J \approx n[R(\xi_{\mathrm{C}} - x)b - \gamma x]$, analogous to equation (10), where R is the molecule formation rate on grains.

Our results on thermal-chemical instabilities are relevant for condensation, fragmentation, and collapse of interstellar clouds. A satisfactory understanding of their role in the evolution of interstellar clouds cannot be achieved, however, without a full hydrodynamic treatment. Nonetheless, some evolutionary possibilities may be inferred from the instability analysis. The high temperature instabilities, associated with optical depth effects, may cause warm diffuse clouds to evolve toward thicker and denser configurations, which may undergo further condensation at low temperatures. The pressure loops in Figure 1 have intersections with lines of constant pressure at two states which would be stable in the absence of optical depth effects. For $\tau_v = 0.75$, for example, these states have (T, n) given by $\sim(25\ \mathrm{K}, 450\ \mathrm{cm}^{-3})$ and $(10\ \mathrm{K}, 1000\ \mathrm{cm}^{-3})$, corresponding to a pressure of $\sim 5 \times 10^3$ K cm^{-3}. The thermal-chemical instabilities can lead to fragmentation of a molecular cloud, because their growth times, in the range from 2×10^5–10^6 yr, are less than or comparable with the Jean's time $\tau_J \approx 3.4 \times 10^7 n^{-1/2}$ yr. In other words, a gravitationally unstable mass M_J can be produced by thermal-chemical instability within a time $\sim 10^6$ yr. The value of M_J decreases from $\sim 40\ M_\odot$ to $\sim 8\ M_\odot$ in going from the low- to the high-density phase of $\tau_v = 0.75$ pressure loop. Because the optical depth also increases in going to the denser phase, the Jean's mass can decrease further. In conclusion, the evolution of a molecular cloud is affected by numerous phenomena which need to be analyzed simultaneously in a more complete hydrodynamical treatment. We suggest that thermal-chemical instabilities of the type discussed here play a role in the fragmentation of clouds and in the formation of protostellar objects.

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